

# Anomalous $gt\bar{t}$ couplings in the Littlest Higgs Model with T-parity

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## Abstract

In this work we calculate the *leading* electroweak (EW) corrections to the anomalous  $gt\bar{t}$  coupling in the Littlest Higgs model with T-parity (LHT), by applying the Goldstone boson equivalence theorem. In the LHT model, such electroweak corrections arise from the loop diagrams of heavy fermions and the “would-be” Goldstone bosons. We further examine the EW corrections in the top quark pair production via the quark annihilation process at the LHC. The negative EW corrections in the Standard Model are partially canceled by the positive EW corrections from the loops of the new heavy particles, and the latter dominates in the large invariant mass of the top quark pair.

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## I. INTRODUCTION

The top quark is a special quark in the Standard Model (SM) due to its large mass. As the top quark mass ( $m_t$ ) is close to the electroweak symmetry breaking (EWSB) scale,  $m_t \sim 170.9$  GeV [1], studying the top quark physics might shed lights on the mechanism of EWSB. At the Tevatron, the top quark pair is mainly produced via the quark-antiquark annihilation, whereas at the CERN Large Hadron Collider (LHC) it is produced mainly through gluon-gluon fusion. The LHC will be a true top factory, producing hundreds of millions of top quarks every year. With such a large rate, it becomes possible to accurately measure the total cross section of the top quark pair production, which provides a good probe of searching for new physics (NP). The NP effects can modify the  $g\bar{t}t$  coupling via quantum corrections. The non-SM one-loop corrections to the top quark pair production at hadron colliders have been studied within the general two-Higgs-doublet model (2HDM) [2, 3, 4, 5] and the minimal supersymmetric Standard Model (MSSM) [4, 5, 6, 7, 8, 9, 10, 11, 12, 13, 14, 15, 16, 17, 18]. Within these corrections, the Yukawa electroweak radiative correction is especially interesting because of the existence of the large enhancement to the Yukawa couplings in the 2HDM [19] and MSSM [20, 21]. Significant effects indeed were found on both total cross section and differential cross section distributions, as compared to the one-loop electroweak corrections in the SM [2, 22, 23, 24, 25, 26, 27, 28]. In this study we shall examine the leading electroweak corrections to the top quark pair production in the Littlest Higgs model with T-parity [29, 30, 31].

In Little Higgs models [32, 33, 34, 35, 36], the electroweak symmetry is collectively broken and a weak scale Higgs boson mass is radiatively generated. At one-loop order, the large quadratically divergent correction to the Higgs boson mass squared induced by the top quark ( $t$ ) is canceled by its fermionic partner, and that induced by the electroweak gauge bosons are canceled by their bosonic partners. Constraints from the low energy precision data, especially the  $\rho$ -parameter measurement, require that the symmetry breaking scale of the Little Higgs models has to be so high that the predicted phenomenology has little relevance to the current high energy collider physics program [37, 38, 39]. To alleviate the constraints from low energy data, a discrete symmetry, called T-parity [29, 30, 31], is introduced and warrants the  $\rho$ -parameter to be one at tree-level. In order to incorporate the T-parity systematically, extra fermion fields have to be introduced. As a result, we have

two sets of particles: the usual SM particles and an additional  $T_+$  quark are “even” under the T-parity while the other heavy new particles are “odd”. The SM gauge bosons do not mix with the heavy gauge bosons due to the T-parity, and the corrections to the low energy observables are loop-suppressed, consequently, the new particle mass scale  $f$  of the model as low as 500 GeV is still allowed [40]. Thus the masses of the new particles are at the order of TeV, and they may cause large quantum corrections to the top quark pair production at high energy colliders. In this paper, we calculate the *leading* electroweak (EW) radiative corrections to the anomalous  $gt\bar{t}$  couplings by applying the Goldstone-boson equivalence theorem (ET) [41, 42, 43, 44, 45, 46, 47, 48, 49, 50, 51, 52, 53, 54, 55, 56]. We also examine their effects in the  $q\bar{q} \rightarrow g \rightarrow t\bar{t}$  processes at the LHC. The one-loop leading EW corrections to the anomalous  $gt\bar{t}$  coupling are given in terms of the Passarino-Veltman scalar functions [57], which are evaluated using the library LOOPTOOLS (FF) [58, 59, 60].

## II. LITTLEST HIGGS MODEL WITH T-PARITY

The Littlest Higgs model with T-parity (LHT) is based on a  $SU(5)/SO(5)$  nonlinear sigma model whose low energy Lagrangian is described in detail in Refs. [29, 30, 31, 61, 62]. With the global symmetry  $SU(5)$  being broken down to  $SO(5)$  by a  $5 \times 5$  symmetric tensor at the scale  $f$ , the gauged  $[SU(2) \times U(1)]_1 \times [SU(2) \times U(1)]_2$ , a subgroup of  $SU(5)$ , is broken to the diagonal  $SU(2)_W \times U(1)_Y$ , a subgroup of  $SO(5)$ . Four new (T-odd) heavy gauge bosons appear after the symmetry breaking: the photon partner ( $A_H$ ), the  $Z$ -boson partner ( $Z_H$ ) and the  $W^\pm$ -boson partner ( $W_H^\pm$ ). We shall apply the ET to calculate the leading electroweak Yukawa contributions and adopt the following notations:  $h$  is the Higgs boson;  $\pi^0(\pi^\pm)$  is the Goldstone-boson (GB) eaten by the  $Z$ -boson ( $W$ -boson);  $\omega^0(\omega^\pm, \eta)$  is the Goldstone-boson eaten by  $Z_H(W_H, A_H)$ <sup>1</sup>. Furthermore, a copy of leptons and quarks with T-odd quantum numbers are added in order to preserve the T-parity. The T-odd heavy quarks which contribute to the  $gt\bar{t}$  coupling are  $t_-$ ,  $b_-$  and  $T_-$ , which are T-parity partners of the SM top, bottom quarks and heavy T-even  $T_+$  quark, respectively. The interactions between the SM top quark, the  $T_+$  quark, scalars (the Higgs boson and GBs), and T-odd

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<sup>1</sup> There is an order of  $v^2/f^2$  mixing between  $\omega^\pm$  and the  $SU(2)$  triplet T-odd scalars  $\phi^\pm$  [40], which is neglected in our calculation.

Table I: The relevant couplings of the SM top quark and new particles.

	$h - t - T_+$	$\pi^0 - t - T_+$	$\omega^0 - t - t_-$	$\eta - t - t_-$	$\omega^- - t - b_-$	$\eta - t - T_-$
$g_V$	$-\frac{\lambda_1^2}{2\sqrt{\lambda_1^2 + \lambda_2^2}}$	$i\frac{\lambda_1^2}{2\sqrt{\lambda_1^2 + \lambda_2^2}}$	$i\frac{\sqrt{2}}{4}\kappa$	$-i\frac{\sqrt{10}}{20}\kappa$	$i\frac{1}{2}\kappa$	$-i\frac{\sqrt{5}}{5}\frac{\lambda_1\lambda_2}{\sqrt{\lambda_1^2 + \lambda_2^2}}$
$g_A$	$-\frac{\lambda_1^2}{2\sqrt{\lambda_1^2 + \lambda_2^2}}$	$i\frac{\lambda_1^2}{2\sqrt{\lambda_1^2 + \lambda_2^2}}$	$i\frac{\sqrt{2}}{4}\kappa$	$-i\frac{\sqrt{10}}{20}\kappa$	$i\frac{1}{2}\kappa$	$i\frac{\sqrt{5}}{5}\frac{\lambda_1\lambda_2}{\sqrt{\lambda_1^2 + \lambda_2^2}}$

quarks could be found by expanding the following effective Lagrangian,

$$\mathcal{L}_t = -\frac{\lambda_1}{2\sqrt{2}}f\epsilon_{ijk}\epsilon_{xy}\left[(\bar{Q}_1)_i\Sigma_{jx}\Sigma_{ky} - (\bar{Q}_2\Sigma_0)_i\tilde{\Sigma}_{jx}\tilde{\Sigma}_{ky}\right]u_R - \lambda_2f(\bar{U}_1U_{R1} + \bar{U}_2U_{R2}) + h.c. \quad (1)$$

and

$$\mathcal{L}_\kappa = -\kappa f[\bar{\Psi}_2\xi\Psi_c + \bar{\Psi}_1\Sigma_0(\Omega\xi^\dagger\Omega)\Psi_c] + h.c., \quad (2)$$

where  $\epsilon_{ijk}$  and  $\epsilon_{xy}$  are antisymmetric tensors, and  $i, j, k$  run over  $1 - 3$  and  $x, y$  over  $4 - 5$ ;  $Q_1 = (q_1, U_1, 0, 0)^T$ ,  $Q_2 = (0, 0, U_2, q_2)^T$  where  $q_i = -\sigma_2(u_i, d_i)^T = (id_i, -iu_i)^T$  with  $i = 1, 2$ ;  $\Psi_i = (q_i, 0, 0, 0)^T$  and  $\Psi_c = (q_c, \chi_c, \tilde{q}_c)^T$ . (Here, the superscript  $T$  denotes taking transpose.) Also,  $\Sigma = \xi^2\Sigma_0$  and  $\tilde{\Sigma} = \Sigma_0\Omega\Sigma^\dagger\Omega\Sigma_0$  which is the T-parity transformation of  $\Sigma$ , where  $\xi = \exp\{i\Pi^a X^a/f\}$ ,  $X^a$  are the broken generators,  $\Pi^a$  contain the Higgs boson and all the other GB fields, and

$$\Sigma_0 = \begin{bmatrix} 0_{2\times 2} & 0_{2\times 1} & 1_{2\times 2} \\ 0_{1\times 2} & 1 & 0_{1\times 2} \\ 1_{2\times 2} & 0_{2\times 1} & 0_{2\times 2} \end{bmatrix} \text{ and } \Omega = \begin{bmatrix} 1_{2\times 2} & & \\ & -1 & \\ & & 1_{2\times 2} \end{bmatrix}_{5\times 5}. \quad (3)$$

For more details of the LHT model, see Refs. [29, 30, 31, 61, 62]. Here, we only list the couplings of the SM top quark and new heavy particles, which contribute to the loop corrections to the  $gt\bar{t}$  coupling, as shown in Table I<sup>2</sup>. The coupling of the  $\bar{t}FS$  interaction relevant to our calculations is given as  $i(g_V + g_A\gamma_5)$ , where  $F$  ( $S$ ) denotes the heavy fermion (scalar). There also exist couplings between T-odd  $SU(2)$  triplet scalars  $\phi$  to the top quark, but they are neglected in this work since they are at the  $\mathcal{O}(v/f)$ . Since we perform our calculations in the 't Hooft-Feynman gauge, the mass of the would-be GB is the same as its corresponding gauge boson. The masses of the heavy particles are given as follows:

$$m_t \sim \frac{\lambda_1\lambda_2}{\sqrt{\lambda_1^2 + \lambda_2^2}}v, \quad m_{T_+} \sim \sqrt{\lambda_1^2 + \lambda_2^2}f, \quad m_{T_-} = \lambda_2f, \quad (4)$$

<sup>2</sup> Our Feynman rules coincide with the results in Refs. [61, 63, 64], up to the  $\mathcal{O}(v/f)$  accuracy.

$$m_{\omega^{\pm,0}} \sim gf, \quad m_\eta \sim \frac{g'f}{\sqrt{5}}, \quad m_{t-} \simeq m_{b-} \sim \sqrt{2}\kappa f, \quad (5)$$

where  $g$  ( $g'$ ) is the weak (hypercharge) gauge coupling strength, and  $v \simeq 246$  GeV. With those couplings and masses of the new particles, we now calculate the one-loop corrections to the  $gt\bar{t}$  coupling in the LHT model.

### III. FORM FACTOR OF $gt\bar{t}$ AND ONE-LOOP EW CORRECTIONS IN THE LHT

Following the parametrization in Ref. [2], the effective matrix element of  $gt\bar{t}$ , including the one-loop corrections, can be written as

$$-ig_s T^a \bar{u}_t \Gamma^\mu v_{\bar{t}}, \quad (6)$$

with

$$\Gamma^\mu = (1 + \alpha)\gamma^\mu + i\beta\sigma^{\mu\nu}q_\nu + \xi \left( \gamma^\mu - \frac{2m_t}{\hat{s}} q^\mu \right) \gamma_5. \quad (7)$$

where the loop-induced form factors  $\alpha$ ,  $\beta$  and  $\xi$  are usually refereed as the chromo-charge, chromo-magnetic-dipole<sup>3</sup> and chromo-anapole, respectively. Here,  $g_s$  is the strong coupling strength,  $T^a$  are the color generators,  $q = p_t + p_{\bar{t}}$ , and  $\hat{s} = q^2$ . After summing over the final state and averaging over the initial state colors and spins, the constituent total cross section of  $q\bar{q} \rightarrow g \rightarrow t\bar{t}$  is [2]

$$\hat{\sigma} = \frac{8\pi\alpha_s^2}{27\hat{s}^2} \sqrt{1 - \frac{4m_t^2}{\hat{s}}} \left\{ \hat{s} + 2m_t^2 + 2\Re \left[ (\hat{s} + 2m_t^2)\alpha + 3m_t\hat{s}\beta \right] \right\}, \quad (8)$$

where  $\alpha_s \equiv g_s^2/(4\pi)$ , and  $\Re$  denotes taking its real part. Note that  $\xi$  does not contribute to the total cross section at this order, as a result of the interference with the Born matrix

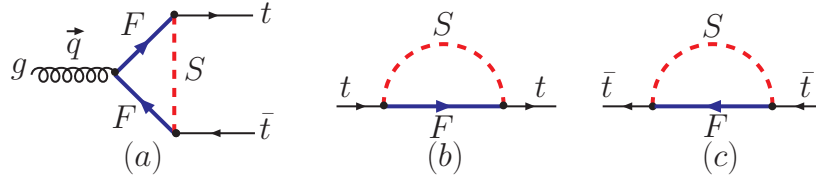


Figure 1: Feynman diagrams of the one-Loop corrections to the  $gt\bar{t}$  coupling in the LHT model .

<sup>3</sup> The one-loop non-SM contributions to the  $gt\bar{t}$  chromo-magnetic-dipole form factor have been recently studied in the literature [65], where several models are considered, including 2HDM, topcolor assisted Technicolor model, 331 model and universal extra dimension model.

element, but for completeness we will present the analytical expressions of those three form factors in the LHT model below.

At the one-loop level, the  $gt\bar{t}$  coupling receives two kinds of quantum corrections: one is the irreducible triangle-loop correction (Fig. 1a), another is the self-energy correction to the external top quark lines (Figs. 1b and 1c). For simplicity, we use the particles running inside the loop to represent the corresponding loop correction diagram. For example, Fig. 1(a) are denoted as  $(F, F, S)$ . In the LHT model, the diagrams contributing to the anomalous  $gt\bar{t}$  coupling are given by  $(T_+, T_+, h/\pi^0)$ ,  $(t_-, t_-, \eta/\omega^0)$ ,  $(b_-, b_-, \omega^\pm)$  and  $(T_-, T_-, \eta)$ . The coupling strength of the  $gF\bar{F}$  vertex is just the usual strong coupling while the  $\bar{t}FS$  couplings are explicitly given in Table I.

We use the dimensional regularization scheme to regulate the ultraviolet divergences and adopt the on-mass-shell renormalization scheme to renormalize the electroweak parameters. In this scheme, the wave function renormalization corrections of the external top quark legs are canceled by the corresponding counterterms. We will regularize the ultraviolet divergences in our calculation by dimensional regularization with the regulator defined by  $\Delta = \frac{1}{\epsilon} - \gamma_E + \ln 4\pi$ , where  $2\epsilon \equiv 4 - n$ ,  $n$  is the dimension of the space-time and  $\gamma_E$  is the Euler constant. As we are calculating the leading EW corrections to the  $gt\bar{t}$  coupling, we do not need to introduce the counterterm for the strong coupling. By introducing appropriate counterterms, one can easily deduce the renormalized vertex of  $gt\bar{t}$  as

$$-ig_s T^a \bar{u}_t (\gamma^\mu + \delta\Gamma_{\text{ren}}^\mu) v_t, \quad (9)$$

where

$$\delta\Gamma_{\text{ren}}^\mu = \gamma^\mu (\delta Z_V^t + \delta Z_A^t \gamma_5) + \delta\Gamma_\Delta^\mu \quad (10)$$

Here,  $\delta Z_{V,A}^t$  denote the wave function renormalization constants of the external top quark lines, which are defined by

$$Z_t \equiv 1 + \delta Z_t = 1 + \delta Z_V^t + \delta Z_A^t \gamma_5,$$

while  $\delta\Gamma_\Delta$  denotes the triangle loop corrections to the vertex. Clearly, the  $\delta Z_V$  counter terms only contribute to the form factor  $\alpha$ , the  $\delta Z_A$  counter terms only contribute to the form factor  $\xi$ , but the vertex corrections  $\delta\Gamma_\Delta$  contribute to all three form factors. We thus write the form factors as follows,

$$\alpha = \alpha_\Delta + \delta Z_V, \quad \beta = \beta_\Delta, \quad \xi = \xi_\Delta + \delta Z_A, \quad (11)$$

where  $\alpha_\Delta$ ,  $\beta_\Delta$  and  $\xi_\Delta$  denote the coefficients of the  $\gamma^\mu$ ,  $\sigma^{\mu\nu}q_\nu$  and  $\gamma^\mu\gamma_5$  terms in  $\delta\Gamma_\Delta^\mu$ , respectively. Note that there is an additional term  $q^\mu\gamma_5$  in  $\delta\Gamma_\Delta^\mu$ . After adding the  $\delta Z_A$  counter terms, we can write the combination of  $\gamma^\mu\gamma_5$  and  $q^\mu\gamma_5$  in a compact form as the  $\xi$  term in Eq. (7), which is guaranteed by the Ward identity for the conservation of QCD current.

Consider the renormalization constants. The wave function renormalization constants can be determined from the top quark self-energy diagrams, cf. Figs. 1(b, c), which can be decomposed as follows:

$$\Sigma(\not{p}) = \not{p} [\Sigma_V(p^2) + \Sigma_A(p^2) \gamma_5] + m_t \Sigma_S(p^2). \quad (12)$$

In the on-shell scheme, the finite parts of the counter terms are determined by the requirement that the residue of the fermion propagator is equal to one, which fixes the wave function renormalization constraints by

$$\delta Z^V = -\Sigma_V(p^2 = m_t^2) - 2m_t^2 \frac{\partial}{\partial p^2} (\Sigma_V + \Sigma_S)|_{p^2=m_t^2}, \quad (13)$$

$$\delta Z^A = -\Sigma_A(p^2 = m_t^2). \quad (14)$$

In the LHT model, they are given by

$$\begin{aligned} \delta Z^V &= \frac{1}{16\pi^2} \frac{g_V^2 + g_A^2}{2m_t^2} \{A_0(m_S^2) - A_0(m_F^2) + (m_F^2 - m_S^2 - m_t^2) B_0(m_t^2)\} \\ &\quad + \frac{1}{16\pi^2} \left[ (g_V^2 + g_A^2) (-m_t^2 + m_S^2 - m_F^2) - (g_V^2 - g_A^2) 2m_t m_F \right] B'_0(m_t^2), \end{aligned} \quad (15)$$

$$\delta Z^A = \frac{1}{16\pi^2} \frac{g_V g_A}{m_t^2} \{A_0(m_S^2) - A_0(m_F^2) + (m_F^2 - m_S^2 + m_t^2) B_0(m_t^2)\}, \quad (16)$$

where  $A_0$  and  $B_0$  are the well-known one-point and two-point scalar functions [57]. We also introduce the following shorthand notations,

$$B_0(m_t^2) \equiv B_0(m_t^2; m_S^2, m_F^2), \quad B'_0(m_t^2) \equiv \frac{\partial}{\partial p^2} B_0(p^2; m_S^2, m_F^2)|_{p^2=m_t^2}. \quad (17)$$

where  $m_S$  ( $m_F$ ) is the mass of the scalar (fermion) in the loop.

Now considering the vertex corrections  $\delta\Gamma_\Delta^\mu$ , which we decompose into the form factors  $\alpha_\Delta$ ,  $\beta_\Delta$  and  $\xi_\Delta$ , as listed below. The form factor  $\alpha_\Delta$  is given by

$$\begin{aligned} \alpha_\Delta &= -\frac{g_V g_V^*}{16\pi^2} \{ \alpha_1 + \alpha_2 B_0(\hat{s}) + \alpha_3 B_0(m_t^2) + \alpha_4 C_0 \} \\ &\quad - \frac{g_A g_A^*}{16\pi^2} \{ \alpha'_1 + \alpha'_2 B_0(\hat{s}) + \alpha'_3 B_0(m_t^2) + \alpha'_4 C_0 \}, \end{aligned} \quad (18)$$

where

$$\alpha_1 = \frac{\hat{s}}{2(\hat{s} - 4m_t^2)} + \frac{2}{\hat{s} - 4m_t^2} [-A_0(m_S^2) + A_0(m_F^2)], \quad (19)$$

$$\alpha_2 = \frac{1}{2(\hat{s} - 4m_t^2)^2} \left[ -16m_t^4 - 32m_F m_t^3 + (-16m_F^2 + 16m_S^2 + 14\hat{s})m_t^2 \right. \\ \left. + 8m_F \hat{s} m_t - \hat{s}^2 - 2m_F^2 \hat{s} + 2m_S^2 \hat{s} \right], \quad (20)$$

$$\alpha_3 = \frac{1}{2(\hat{s} - 4m_t^2)^2} \left[ 32m_F m_t^3 + (32m_F^2 - 32m_S^2 - 6\hat{s})m_t^2 \right. \\ \left. - 8m_F \hat{s} m_t - 2\hat{s}(m_F^2 - m_S^2) \right], \quad (21)$$

$$\alpha_4 = \frac{1}{2(\hat{s} - 4m_t^2)^2} \left[ 16m_t^6 + 32m_F m_t^5 + (32m_F^2 - 32m_S^2 - 6\hat{s})m_t^4 \right. \\ + (32m_F^3 - 32m_F m_S^2 - 24m_F \hat{s})m_t^3 \\ + (16m_F^4 + 16m_S^4 - 32m_F^2 m_S^2 + 2\hat{s}^2 - 28m_F^2 \hat{s} + 20m_S^2 \hat{s})m_t^2 \\ \left. + (4m_F \hat{s}^2 - 8m_F^3 \hat{s} + 8m_F m_S^2 \hat{s})m_t + 2m_F^2 \hat{s}^2 + 2m_F^4 \hat{s} + 2m_S^4 \hat{s} - 4m_F^2 m_S^2 \hat{s} \right], \quad (22)$$

and

$$\alpha'_1 = \alpha_1, \quad \alpha'_{2,3,4} = \alpha_{2,3,4} \Big|_{m_F \rightarrow -m_F}. \quad (23)$$

Here we introduce the following shorthand notations,

$$B_0(\hat{s}) \equiv B_0(\hat{s}; m_t^2, m_t^2), \quad C_0 \equiv C_0(m_t^2, \hat{s}; m_S^2, m_F^2, m_F^2), \quad (24)$$

where  $C_0(\dots)$  is the usual three-point scalar function [57]. The form factor  $\beta_\Delta$  is given by

$$\beta_\Delta = \frac{g_V g_V^*}{16\pi^2} \{ \beta_1 + \beta_2 B_0(\hat{s}) + \beta_3 B_0(m_t^2) + \beta_4 C_0 \} \\ + \frac{g_A g_A^*}{16\pi^2} \{ \beta'_1 + \beta'_2 B_0(\hat{s}) + \beta'_3 B_0(m_t^2) + \beta'_4 C_0 \}, \quad (25)$$

where

$$\beta_1 = \frac{m_t}{\hat{s} - 4m_t^2} + \frac{1}{m_t(\hat{s} - 4m_t^2)} [-A_0(m_S^2) + A_0(m_F^2)], \quad (26)$$

$$\beta_2 = \frac{1}{(\hat{s} - 4m_t^2)^2} [2m_t^3 - 8m_F m_t^2 + (-6m_F^2 + 6m_S^2 + \hat{s})m_t + 2m_F \hat{s}], \quad (27)$$

$$\beta_3 = \frac{1}{m_t(\hat{s} - 4m_t^2)^2} \left[ -2m_t^4 + 8m_F m_t^3 + (10m_F^2 - 10m_S^2 - \hat{s})m_t^2 \right. \\ \left. - 2m_F \hat{s} m_t + (m_S^2 - m_F^2)\hat{s} \right], \quad (28)$$

$$\beta_4 = \frac{-2}{(\hat{s} - 4m_t^2)^2} \left[ m_t^5 + 4m_F m_t^4 + (2m_F^2 + 2m_S^2 - \hat{s})m_t^3 - m_F(4m_F^2 - 4m_S^2 + \hat{s})m_t^2 \right. \\ \left. + (-3m_F^4 + m_F^2(6m_S^2 + \hat{s}) - 3m_S^4 - 2m_S^2 \hat{s})m_t + m_F(m_F^2 - m_S^2)\hat{s} \right], \quad (29)$$



and

$$\beta'_1 = \beta_1, \quad \beta'_{2,3,4} = \beta_{2,3,4} \Big|_{m_F \rightarrow -m_F}. \quad (30)$$

Finally, the form factor  $\xi_\Delta$  is given by

$$\xi_\Delta = -\frac{g_V g_A^*}{16\pi^2} \{ -1 + \xi_1 B_0(\hat{s}) + \xi_2 B_0(m_t^2) + \xi_3 C_0 \}, \quad (31)$$

where

$$\xi_1 = \frac{1}{\hat{s} - 4m_t^2} [2m_t^2 - 2m_S^2 + 2m_F^2 + \hat{s}], \quad (32)$$

$$\xi_2 = \frac{-2}{\hat{s} - 4m_t^2} [m_F^2 - m_S^2 + 3m_t^2], \quad (33)$$

$$\xi_3 = \frac{-2}{\hat{s} - 4m_t^2} [m_t^4 - (2m_F^2 + 2m_S^2 + \hat{s})m_t^2 + m_S^4 + m_F^4 - 2m_F^2 m_S^2 + m_F^2 \hat{s}]. \quad (34)$$

#### IV. NUMERICAL RESULTS

The model parameters for the numerical evaluation are  $\lambda_1$ ,  $\lambda_2$ ,  $\kappa$  and  $f$ . As  $\lambda_1$  and  $\lambda_2$  are related by the mass of the top quark, cf. Eq. (4), we could choose either one as the input parameter, and in this study  $\lambda_1$  is chosen. As pointed out from the partial wave study in Ref. [63],  $\lambda_1$  should be bounded in the region  $0.71 \lesssim \lambda_1 \lesssim 2.51$ . Furthermore, if  $\kappa$  is not universal for quark and lepton sectors, as studied in Ref. [66], the upper bound for  $\kappa$  of the quark sector from the constraints of four-fermion operators could be quite loose even for a low  $f$  value, say  $f \sim 500$  GeV. For illustration, we choose the values of the parameters as follows:

$$\begin{aligned} \lambda_1 &= 2.5, \quad \kappa = 5, \quad f = 500 \text{ GeV}, \quad m_t = 175 \text{ GeV}, \\ m_W &= 80.4 \text{ GeV}, \quad m_Z = 91.2 \text{ GeV}, \quad m_h = 120(500) \text{ GeV}, \end{aligned}$$

where  $m_W$ ,  $m_Z$  and  $m_h$  denote the masses of the  $W$ -boson,  $Z$ -boson and Higgs boson, respectively, and the bottom quark is considered as massless throughout this work. With the chosen parameters, the masses of new heavy particles are given by

$$\begin{aligned} m_{T_+} &= 1302 \text{ GeV}, \quad m_{T_-} = 364 \text{ GeV}, \\ m_{t_-} &\simeq m_{b_-} = 3536 \text{ GeV}, \quad m_{\omega^{\pm,0}} = 327 \text{ GeV}, \quad m_\eta = 78 \text{ GeV}. \end{aligned}$$

Since, as a result of the interference with the Born matrix element,  $\xi$  does not contribute, we need only the form factors  $\alpha$  and  $\beta$ , which depend on both the couplings ( $g_V$  and  $g_A$ )

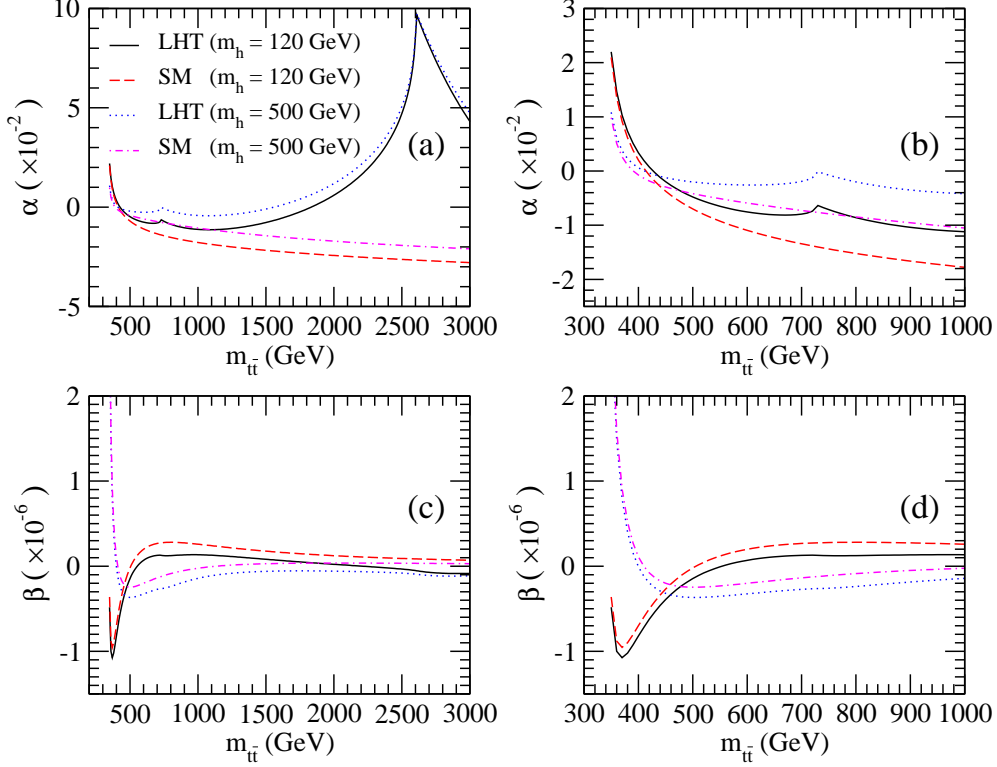


Figure 2: Dependence of the invariant mass of the top quark pair in form factors in both the LHT model and the SM: (a) and (b)  $\alpha$ ; (c) and (d)  $\beta$ . (b) and (d) is the same as (a) and (c), respectively, but focusing on the small  $m_{t\bar{t}}$  region.

and the masses of the scalars and fermions flowing in the loops. We split the form factors in the LHT,  $\alpha_{LHT}$  and  $\beta_{LHT}$ , as follows:

$$\alpha_{LHT} = \alpha_{SM} + \alpha_{HEAVY}, \quad \beta_{LHT} = \beta_{SM} + \beta_{HEAVY}, \quad (35)$$

where the subscript *SM* and *HEAVY* denote contributions to form factors which are induced by the SM loops and the new heavy particle loops, respectively. In Figs. 2(a) and (c), we present the values of form factors  $\alpha$  and  $\beta$  as a function of the invariant mass of the top quark pair system, respectively. In order to investigate the dependence of the SM Higgs boson mass, we also choose two different Higgs boson masses:  $m_h = 120$  GeV and  $m_h = 500$  GeV. We note a few interesting points listed as follows:

- For  $m_{t\bar{t}} > 500$  GeV,  $\alpha_{SM}$  is negative but  $\alpha_{HEAVY}$  is positive. Furthermore, in the region of  $400 \text{ GeV} < m_{t\bar{t}} < 2000 \text{ GeV}$ ,  $\alpha_{HEAVY} \simeq |\alpha_{SM}|$ . Therefore, their sum,  $\alpha_{LHT}$ , is around zero. The small kink in  $\alpha_{HEAVY}$  near  $m_{t\bar{t}} \sim 2m_{T_-}$  GeV is due to the threshold

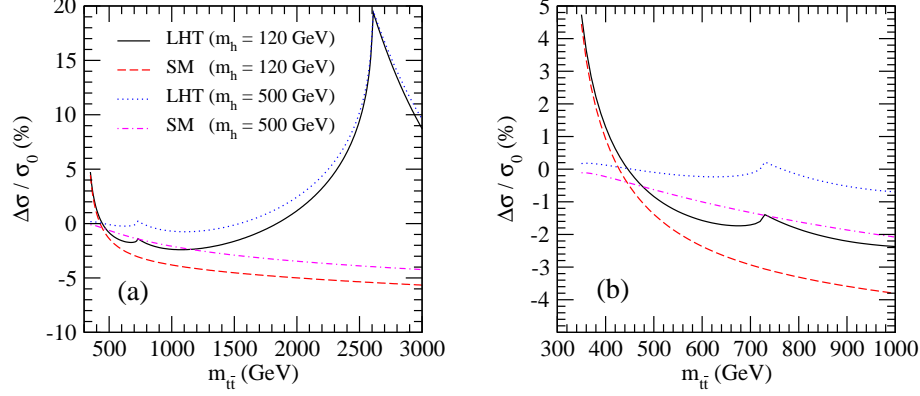


Figure 3: The ratio of the one-loop leading EW correction to the Born level total cross section of  $q\bar{q} \rightarrow g \rightarrow t\bar{t}$  at the LHC. (b) is the same as (a) but focusing on the small  $m_{t\bar{t}}$  region.

effect from producing the  $T_- \bar{T}_-$  pair. However, in the large  $m_{t\bar{t}}$  region, e.g.  $m_{t\bar{t}} > 2500$  GeV,  $\alpha_{HEAVY}$  receives a large corrections from the  $(T_+, T_+, h/\pi^0)$  loops, and is much larger than  $|\alpha_{SM}|$ . In particular,  $\alpha_{HEAVY}$  reaches its maximum around the threshold region, i.e.  $m_{t\bar{t}} \sim 2m_{T_+}$ . As a result,  $\alpha_{LHT}$  is positive and much larger than  $\alpha_{SM}$  in the large  $m_{t\bar{t}}$  region, see the (black) solid line ( $m_h = 120$  GeV) and the (blue) dotted line ( $m_h = 500$  GeV) in Fig. 2 (a). In the small  $m_{t\bar{t}}$  region, i.e.  $m_{t\bar{t}} < 500$  GeV,  $\alpha_{HEAVY}$  is negligible and  $\alpha_{LHT} \simeq \alpha_{SM}$ .

- The form factor  $\beta_{HEAVY}$  is always negative, see the (black) solid line (LHT) and the (red) dashed line (SM) in Fig. 2(d). In the large  $m_{t\bar{t}}$  region, both  $\beta_{LHT}$  and  $\beta_{SM}$  are negligible. Note that the chromo-magnetic-dipole form factor  $\beta$  can contribute to the branching ratio of  $b \rightarrow s\gamma$  process [67, 68, 69], and our numerical results are consistent with the current bounds [69].

Below, we will examine the effects of the leading EW corrections on the top quark pair production at the LHC. For that, we calculate the differential cross section,  $d\sigma/dm_{t\bar{t}}$ , given by

$$\frac{d\sigma}{dm_{t\bar{t}}} = \int dx_1 dx_2 \left\{ f_{q/p}(x_1, Q) f_{\bar{q}/p}(x_2, Q) \frac{d\hat{\sigma}}{dm_{t\bar{t}}}(q\bar{q} \rightarrow t\bar{t}) + (x_1 \leftrightarrow x_2) \right\},$$

where  $\hat{\sigma}$  labels the hard process cross section, and  $f_{q/p}(x, Q)$  denotes the parton distribution function of finding the parton  $q$  in the colliding proton with the momentum fraction  $x$ .  $Q$  is the factorization scale of the hard scattering process. In our calculations, we use the CTEQ 6.1 parton distribution functions [70]. We note that at the LHC, the dominant mechanism

for top quark pair production is via gluon-gluon fusion, i.e.,  $gg \rightarrow t\bar{t}$ . Nevertheless, in this work, we focus on the new physics effect predicted by the LHT to top quark pair production cross section in the quark and anti-quark scattering processes. To examine in detail the effect of leading EW corrections, we calculate the relative corrections defined as

$$\frac{\Delta\sigma}{\sigma^0} \equiv \left( \frac{d\sigma}{dm_{t\bar{t}}} - \frac{d\sigma_0}{dm_{t\bar{t}}} \right) / \frac{d\sigma_0}{dm_{t\bar{t}}}, \quad (36)$$

where  $\sigma_0$  denotes the tree-level SM cross section. Fig. 3(a) shows our numerical results, while Fig. 3(b) reveals the details of the small  $m_{t\bar{t}}$  region of Fig. 3(a). It is clear that the relative corrections are dominated by  $\alpha$ , because  $\alpha$  is much larger than  $\beta$ . Again, we find that the negative EW corrections in the SM are almost canceled by the positive EW corrections from the new heavy particle loops in the LHT model in the region of  $m_{t\bar{t}} < 2000 \text{ GeV}$ . In the large  $m_{t\bar{t}}$  region, the leading EW corrections in the LHT model could increase the cross section by about 20%. However, such a deviation might hardly be recognized as the cross section drops rapidly with increasing  $m_{t\bar{t}}$ . Moreover, bearing in mind that the top quark pair production at the LHC is predominately via the gluon-gluon fusion process, a systematic study including the  $gg \rightarrow t\bar{t}$  process is in order and will be presented in the forthcoming paper.

## V. CONCLUSION

In this paper, we calculate the leading electroweak (EW) corrections to the anomalous  $gt\bar{t}$  couplings in the LHT model by applying the Goldstone-boson equivalence theorem, and further examine their effects on the top quark pair production cross section via quark annihilation processes at the LHC. We found that the negative EW corrections in the SM are partially canceled by the positive EW corrections from the new heavy particle loops in the LHT model. The net one-loop electroweak correction is close to zero in the range of  $500 \text{ GeV} < m_{t\bar{t}} < 2000 \text{ GeV}$ . For a larger value of  $m_{t\bar{t}}$ , the new heavy particle loop correction dominates. A complete study including the electroweak corrections to the top quark pair production via the gluon-gluon fusion process will be presented in the forthcoming paper.

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